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Spin current in ferromagnet/insulator/superconductor junctions

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Abstract

A theory of spin polarized tunneling spectroscopy based on a scattering theory is given for tunneling junctions between ferromagnets and d -wave superconductors. The spin filtering effect of an exchange field in the insulator is also treated. We clarify that the properties of the Andreev reflection are largely modified due to a presence of an exchange field in the ferromagnets, and consequently the Andreev reflected quasiparticle shows an evanescent-wave behavior depending on the injection angle of the quasiparticle. Conductance formulas for the spin current as well as the charge current are given as a function of the applied voltage and the spin-polarization in the ferromagnet for arbitrary barrier heights. It is shown that the surface bound states do not contribute to the spin current and that the zero-bias conductance peak expected for a d -wave superconductor splits into two peaks under the influence of the exchange interaction in the insulator.

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I. INTRODUCTION

The transport properties in hybrid structures between ferromagnets and superconductors have received considerable theoretical and experimental attentions. Interest in such structures includes spin-dependent spectroscopy of superconductors and possible device applications. Since the Cooper pairs in spin singlet superconductors are formed between up and down spins, the high density of spin injection through a tunneling barrier induces a spin imbalance. This non-equilibrium state is expected to result in a suppression of the critical temperature and the critical current density in the superconductor. A large number of experimental studies on spin-polarized tunneling have already been performed using conventional metal superconductors such as Al and Nb about 20 years ago¹. However, the recent discovery of so-called colossal magneto-resistance (CMR) in Mn oxides compound has aroused new interest in this field^{2,3}, because hybrid structure fabrication of the spin-polarized ferromagnets with high- T_c superconductors is now possible using these materials^{4,5}.

On the other hand, the properties of ferromagnet/insulator/superconductor (F/I/S) and ferromagnet/ferromagnetic-insulator/superconductor (F/FI/S) junctions have been analyzed based on the assumption that the conductance spectra correspond to the density of states (DOS) of the superconductor weighted by the spin polarization^{1,6,7}. A theory for F/I/S junctions based on a scattering method has been presented by de Jong and Beeneker⁸, and new aspects of Andreev reflection have been revealed, and also detailed comparisons between theory and experiments have been accomplished^{9,10}. However, these results are restricted to isotropic s -wave superconductors.

In contrast to s -wave superconductor cases, at the interface of a $d_{x^2-y^2}$ -wave superconductor, zero-energy states (ZES) are formed due to the interference effect of the internal phase of the pair potential¹¹. Tunneling theory for $d_{x^2-y^2}$ -wave superconductors has already been presented by extending the BTK formula¹² to include the anisotropy of the pair potential¹³⁻¹⁵. The theory predicts the existence of zero-bias conductance peak (ZBCP) which reflects the formation of the surface bound states on the d -wave superconductors. In this paper, an exchange interaction is introduced on the normal side of the junction and on the insulator in order to analyze the spin polarized tunneling effects. The bound-state condition and tunneling spectroscopy of ferromagnet/ d -wave superconductor junctions have already been analyzed in two papers^{16,17}. They have revealed several important features in charge transport. Here we will argue that the properties of the Andreev reflection¹⁸ is largely modified due to the presence of the exchange interaction. In particular, the existence of an evanescent type of the Andreev reflection, which is referred to as virtual Andreev reflection (VAR), is explained for the first time (see Ref. 27). This process has significant roles on the transports especially for junctions between half-metallic ferromagnets and superconductors. The conductance formulas for the charge and the spin currents are presented based on the scattering method by fully taking account of the VAR process. The merit of a formula based on the scattering methods is that the conductance spectra can easily be calculated for arbitrary barrier heights cases without the restriction of the high-barrier limit. The spin current is, we believe, the most important physical quantity in spin injection devices based on the following two reasons: one is that the spin current gives a direct criteria to estimate the effect of the spin imbalance induced by the tunneling current, the other is that the charge and the spin conductivity may illuminate the study of electron systems that undergo

spin-charge separation, such as Tomonaga-Luttinger liquids and possibly underdoped high- T_c superconductors^{19–21}. We will also analyze ferromagnetic insulator effects, which includes the spin-filtering effect^{6,7}, due to the presence of an exchange field in the insulator. It is shown that a spin-dependent energy shift during the tunneling process induces a splitting of the ZBCP. Based on the detailed analysis of the conductance spectra, we propose a simple method to distinguish the broken time-reversal symmetry (BTRS) states inducement at the surface^{22–24} from spin-dependent tunneling effects. The implications of the ferromagnetic insulator effects on tunneling experiments of high- T_c superconductors and a proposal for possible device applications are also presented.

II. FORMULATION

For the model of formulation, a planar $F/FI/S$ junction with semi-infinite electrodes in the clean limit is assumed. A flat interface is assumed to be located at $x = 0$, and the insulator for up [down] spin is described by a potential $V_{\uparrow[\downarrow]}(\mathbf{x})$ $\{V_{\uparrow[\downarrow]}(\mathbf{x}) = (\hat{V}_0 - [+]\hat{U}_B)\delta(x)\}$, where $\delta(x)$, \hat{V}_0 and \hat{U}_B are the δ -function, a genuine barrier amplitude and an exchange amplitude in the barrier, respectively. The effective mass m in the ferromagnet and in the superconductor are assumed to be equal. For the model of the ferromagnet, we adopt the Stoner model where the effect of the spin polarization is described by the one-electron Hamiltonian with an exchange interaction similarly to the case of Ref.^{8,16,17}. For the description of the $d_{x^2-y^2}$ -wave superconductor, we apply the quasi-classical approximation where the Fermi energy E_F in the superconductor is much larger than the pair potential following the model by Bruder^{25,26}. The effective Hamiltonian (Bogoliubov-de Gennes equation) is given by

$$\begin{bmatrix} H_0(\mathbf{x}) - \rho U(\mathbf{x}) & \Delta(\mathbf{x}, \theta) \\ \Delta^*(\mathbf{x}, \theta) & -\{H_0(\mathbf{x}) + \rho U(\mathbf{x})\} \end{bmatrix} \begin{bmatrix} u(\mathbf{x}, \theta) \\ v(\mathbf{x}, \theta) \end{bmatrix} = E \begin{bmatrix} u(\mathbf{x}, \theta) \\ v(\mathbf{x}, \theta) \end{bmatrix} \quad (1)$$

Here, E is the energy of the quasiparticle, $U(\mathbf{x})$ is the exchange potential given by $U\Theta(-x)$ ($U \geq 0$) where $\Theta(x)$ is the Heaviside step function, ρ is 1 [-1] for up [down] spins, $\Delta(\mathbf{x}, \theta)$ is the pair potential and $H_0(\mathbf{x}) \equiv -\hbar^2 \nabla^2 / 2m + V(\mathbf{x}) - E_F$. To describe the Fermi surface difference in F and S, we assume $E_F = E_{FN}$ for $x < 0$ and $E_F = E_{FS}$ for $x > 0$. The pair potential $\Delta(\mathbf{x}, \theta)$ is taken as $\Delta(\theta)\Theta(x)$ for simplicity. The number of up [down] spin electrons is described by N_{\uparrow} [N_{\downarrow}]. The polarization and the wave-vector of quasiparticles in the ferromagnet for up [down] spin are expressed as $P_{\uparrow} \equiv \frac{N_{\uparrow}}{N_{\uparrow}+N_{\downarrow}} = \frac{E_{FN}+U}{2E_{FN}}$ [$P_{\downarrow} \equiv \frac{N_{\downarrow}}{N_{\uparrow}+N_{\downarrow}} = \frac{E_{FN}-U}{2E_{FN}}$] and $k_{N,\uparrow} = |\mathbf{k}_{N,\uparrow}| \equiv \sqrt{\frac{2m}{\hbar^2}(E_{FN} + U)}$ [$k_{N,\downarrow} = |\mathbf{k}_{N,\downarrow}| \equiv \sqrt{\frac{2m}{\hbar^2}(E_{FN} - U)}$], respectively⁸.

We assume the quasiparticle injection of up spin electrons at an angle θ_N to the interface normal as shown in Fig. 1. Four possible trajectories exist; they are Andreev reflection (AR), normal reflection (NR), transmission to superconductor as electron-like quasiparticles (ELQ), and transmission as hole-like quasiparticles (HLQ). The spin direction is conserved for NR but not for AR. When the superconductor has $d_{x^2-y^2}$ -wave symmetry, the effective pair potentials for ELQ and HLQ are given by $\Delta_+ \equiv \Delta_0 \cos 2(\theta_S - \beta)$ and $\Delta_- \equiv \Delta_0 \cos 2(\theta_S + \beta)$, respectively, where β is the angle between a -axis of the crystal and the interface normal. Results for various pairing symmetries are obtained by setting proper values to Δ_+ and Δ_- similarly to the previous formulas^{13,15}. The wave vectors of ELQ and HLQ are approximated

by $k_S = |\mathbf{k}_S| \approx \sqrt{\frac{2mE_{FS}}{\hbar^2}}$ following the model by Andreev¹⁸. Since translational symmetry holds along the y -axis direction, the momentum components of all trajectories are conserved ($k_{N,\uparrow} \sin \theta_N = k_{N,\downarrow} \sin \theta_A = k_S \sin \theta_S$)²⁷. Note that θ_N is not equal to θ_A except when $U = 0$, which means retro-reflectivity of AR is broken. Such novel behavior is a consequence of the fact that in the presence of an exchange field the BCS paring is formed not strictly between states of equal but opposite k -vectors, the so-called Fulde-Ferrell effect²⁸. The wave-function in the ferromagnet ($x < 0$) for up [down] spin with injection angle θ_N is described by

$$\begin{pmatrix} u(\mathbf{x}, \theta_N) \\ v(\mathbf{x}, \theta_N) \end{pmatrix} = e^{i\mathbf{k}_{N,\uparrow[\downarrow]}\mathbf{x}} \begin{pmatrix} 1 \\ 0 \end{pmatrix} + a_{\uparrow[\downarrow]}(E, \theta_N) e^{i\mathbf{k}''_{N,\downarrow[\uparrow]}\mathbf{x}} \begin{pmatrix} 0 \\ 1 \end{pmatrix} + b_{\uparrow[\downarrow]}(E, \theta_N) e^{i\mathbf{k}'_{N,\uparrow[\downarrow]}\mathbf{x}} \begin{pmatrix} 1 \\ 0 \end{pmatrix}, \quad (2)$$

where the signs of the x -components of $\mathbf{k}_{N,\uparrow[\downarrow]}$ and $\mathbf{k}'_{N,\uparrow[\downarrow]}$ are the reversed of each other. The reflection probabilities of the two processes are obtained by solving Eq. (1) and by connecting the wave-function and its derivative at $x = 0$.

Next, we will simply explain the Fermi surface effect by assuming up spin injection. Various kinds of reflection process are expected depending on the values of E_{FN} , E_S and U . For example, when $k_S < k_{N,\uparrow}$, total reflection ($|b_{\uparrow[\downarrow]}(E, \theta_N)|^2 = 1$) occurs when $\theta_N > \sin^{-1}(k_S/k_{N,\uparrow}) \equiv \theta_{c1}$ ^{15,17}. In this case, the net currents of the spin and the charge from the ferromagnet to the superconductor vanish. On the other hand, when $k_{N,\downarrow} < k_S < k_{N,\uparrow}$, the x -component of wave-vector in AR process ($\sqrt{k_{N,\downarrow}^2 - k_S^2 \sin^2 \theta_S}$) becomes purely imaginary for $\theta_{c1} > \theta_N > \sin^{-1}(k_{N,\downarrow}/k_{N,\uparrow}) \equiv \theta_{c2}$. In this case, although transmitted quasiparticles from ferromagnet to superconductor do propagate, the Andreev reflected quasiparticles do not propagate (VAR process). A finite amplitude of the evanescent AR process still exists ($|a_{\uparrow}(E, \theta_N)|^2 > 0$) and the net currents of the spin and the charge from the ferromagnet to the superconductor do not vanish. It is easy to check the conservation laws for the charge, the excitation, and the spin on the VAR process following the method presented in Ref.¹². The existence of the VAR process has not been treated in the one-dimensional model⁸ because it is a peculiar feature of a two or three dimensional F/S interface.

The conductance of the junctions are obtained by extending previous formula to include the effect of spin^{12,8,17}. In the following, consider a situation where $k_{N,\downarrow} < k_S < k_{N,\uparrow}$. To analyze the transport properties of an $F/I/S$ junction, two kinds of conductance spectrum are introduced. The conductance for the charge current is defined by the charge flow induced by the up [down] spin quasiparticle injection and is given by

$$\hat{\sigma}_{q,\uparrow[\downarrow]}(E, \theta_N) \equiv \text{Re} \left[1 + \frac{\lambda_2}{\lambda_1} |a_{\uparrow[\downarrow]}(E, \theta_N)|^2 - |b_{\uparrow[\downarrow]}(E, \theta_N)|^2 \right] \quad (3)$$

(for $0 < |\theta_N| < \theta_{c2}$)

$$= \frac{4\lambda_1 [4\lambda_2 |\hat{\Gamma}_+|^2 + (1 + \lambda_2)^2 + Z_{\downarrow[\uparrow]}^2 - |\hat{\Gamma}_+ \hat{\Gamma}_-|^2 \{(1 - \lambda_2)^2 + Z_{\downarrow[\uparrow]}^2\}]}{|(1 + \lambda_1 + iZ_{\uparrow[\downarrow]})(1 + \lambda_2 - iZ_{\downarrow[\uparrow]}) - (1 - \lambda_1 - iZ_{\uparrow[\downarrow]})(1 - \lambda_2 + iZ_{\downarrow[\uparrow]})\hat{\Gamma}_+ \hat{\Gamma}_-|^2}, \quad (4)$$

(for $\theta_{c2} < |\theta_N| < \theta_{c1}$)

$$= \frac{4\lambda_1 (1 - |\hat{\Gamma}_+ \hat{\Gamma}_-|^2) \{1 + (-\kappa_2 + Z_{\downarrow})^2\}}{|(1 + \lambda_1 + iZ_{\uparrow})\{1 - i(\kappa_2 + Z_{\downarrow})\} - (1 - \lambda_1 - iZ_{\uparrow})\{1 + i(\kappa_2 + iZ_{\downarrow})\}\hat{\Gamma}_+ \hat{\Gamma}_-|^2}, \quad (5)$$

(for $\theta_{c1} < |\theta_N| < \pi/2$)
 $=0$,
 where

$$Z_{\uparrow[\downarrow]} = \frac{Z_{0,\uparrow[\downarrow]}}{\cos \theta_S}, \quad Z_{0,\uparrow[\downarrow]} = \frac{2m(\hat{V}_0 - [+]\hat{U}_B)}{\hbar^2 k_S},$$

$$\hat{\Gamma}_{\pm} = \Gamma_{\pm} \exp(\mp i\phi_{\pm}), \quad \exp i\phi_{\pm} = \frac{\Delta_{\pm}}{|\Delta_{\pm}|}, \quad \Gamma_{\pm} = \frac{E - \sqrt{E^2 - |\Delta_{\pm}|^2}}{|\Delta_{\pm}|},$$

$$\lambda_1 = \frac{k_{N,\uparrow[\downarrow]} \cos \theta_N}{k_S \cos \theta_S}, \quad \lambda_2 = \frac{k_{N,\downarrow[\uparrow]} \cos \theta_A}{k_S \cos \theta_S}, \quad \kappa_2 = i\lambda_2 = \frac{\sqrt{k_S^2 \sin^2 \theta_S - k_{N,\downarrow}^2}}{k_S \cos \theta_S}.$$

The conductance for the spin current is defined by the spin imbalance induced by the up [down] spin quasiparticle injection,

$$\hat{\sigma}_{s,\uparrow[\downarrow]}(E, \theta_N) \equiv \text{Re} \left[1 - \frac{\lambda_2}{\lambda_1} |a_{\uparrow[\downarrow]}(E, \theta_N)|^2 - |b_{\uparrow[\downarrow]}(E, \theta_N)|^2 \right] \quad (6)$$

(for $0 < |\theta_N| < \theta_{c2}$)

$$= \frac{4\lambda_1 \left[-4\lambda_2 |\Gamma_+|^2 + (1 + \lambda_2)^2 + Z_{\downarrow[\uparrow]}^2 - |\hat{\Gamma}_+ \hat{\Gamma}_-|^2 \{ (1 - \lambda_2)^2 + Z_{\downarrow[\uparrow]}^2 \} \right]}{|(1 + \lambda_1 + iZ_{\uparrow[\downarrow]})(1 + \lambda_2 - iZ_{\downarrow[\uparrow]}) - (1 - \lambda_1 - iZ_{\uparrow[\downarrow]})(1 - \lambda_2 + iZ_{\downarrow[\uparrow]})\hat{\Gamma}_+ \hat{\Gamma}_-|^2}, \quad (7)$$

(for $\theta_{c2} < |\theta_N| < \theta_{c1}$)

$$= \frac{4\lambda_1(1 - |\hat{\Gamma}_+ \hat{\Gamma}_-|^2) \{1 + (-\kappa_2 + Z_{\downarrow})^2\}}{|(1 + \lambda_1 + iZ_{\uparrow})\{1 - i(\kappa_2 + Z_{\downarrow})\} - (1 - \lambda_1 - iZ_{\uparrow})\{1 + i(\kappa_2 + iZ_{\downarrow})\}\hat{\Gamma}_+ \hat{\Gamma}_-|^2}, \quad (8)$$

(for $\theta_{c1} < |\theta_N| < \pi/2$)
 $=0$.

The Andreev reflected quasiparticles positively contribute to the charge current, but since their spins are reversed, they have negative contribution to the spin current. Second terms in r.h.s. of Eqs. (3) and (6) do not have finite contribution on net current in the VAR process, since the corresponding λ_2 is purely imaginary. The normalized total conductance spectra for the charge current $\sigma_q(E)$ and the spin current $\sigma_s(E)$ are given by

$$\sigma_q(E) = \sigma_{q,\uparrow}(E) + \sigma_{q,\downarrow}(E), \quad (9)$$

$$\sigma_{q,\uparrow[\downarrow]}(E) = \frac{1}{R_N} \int_{-\pi/2}^{\pi/2} d\theta_N \cos \theta_N \hat{\sigma}_{q,\uparrow[\downarrow]}(E, \theta_N) P_{\uparrow[\downarrow]} k_{F,\uparrow[\downarrow]}, \quad (10)$$

$$\sigma_s(E) = \sigma_{s,\uparrow}(E) - \sigma_{s,\downarrow}(E), \quad (11)$$

$$\sigma_s(E) = \frac{1}{R_N} \int_{-\pi/2}^{\pi/2} d\theta_N \cos \theta_N \hat{\sigma}_{s,\uparrow[\downarrow]}(E, \theta_N) P_{\uparrow[\downarrow]} k_{F,\uparrow[\downarrow]}, \quad (12)$$

where

$$R_N = \int_{-\pi/2}^{\pi/2} d\theta_N \cos \theta_N [\hat{\sigma}_{N,\uparrow}(\theta_N) P_{\uparrow} k_{F,\uparrow} + \hat{\sigma}_{N,\downarrow}(\theta_N) P_{\downarrow} k_{F,\downarrow}], \quad (13)$$

$$\hat{\sigma}_{N,\uparrow[\downarrow]}(\theta_N) = \frac{4\lambda_1}{|1 + \lambda_1 + iZ_{\uparrow[\downarrow]}|^2}.$$

In the above, R_N , $\sigma_{q,\uparrow[\downarrow]}(E)$ and $\sigma_{s,\uparrow[\downarrow]}(E)$ correspond to the conductance when the superconductor is in the normal state and the spin-resolved normalized conductance spectra for charge and spin, respectively. The net polarization $J_p(eV)$ as a function of the bias voltage V is give by

$$J_p(eV) = \frac{\int_{-\infty}^{\infty} dE \sigma_s(E) \{f(E - eV) - f(E)\}}{\int_{-\infty}^{\infty} dE \sigma_q(E) \{f(E - eV) - f(E)\}}, \quad (14)$$

where $f(E)$ is the Fermi distribution function. Since the convolution with $f(E)$ gives only a smearing effect in the conductance spectra, the temperature is set to zero in the following discussions.

In the above formulation, we have neglected the self-consistency of the pair potential in order to get analytical formulas²⁹. However, the present formula is easily extended to include this effect simply by replacing Γ_{\pm} with $\Gamma_{\pm}(x)|_{x=0}$, where $\Gamma_{\pm}(x)$ follows the Ricatti equations described by

$$\frac{d}{dx} \hat{\Gamma}_+(x) = \frac{1}{i\hbar^2 k_F \cos \theta_S} [-\Delta_+(x) \hat{\Gamma}_+^2(x) - \Delta_+^*(x) + 2E \hat{\Gamma}_+(x)], \quad (15)$$

$$\frac{d}{dx} \hat{\Gamma}_-(x) = \frac{1}{i\hbar^2 k_F \cos \theta_S} [-\Delta_-^*(x) \hat{\Gamma}_-^2(x) - \Delta_-(x) + 2E \hat{\Gamma}_-(x)], \quad (16)$$

Here the spatial dependence of the pair potential is assumed as $\Delta_{\pm}(x)$ (functions of x).

The most important differences in the present formula from previous ones are; i) a novel formula for the non-linear spin current, ii) a capability to treat the ferromagnetic insulator effects based on the scattering method, iii) the introduction of the breakdown in the retro-reflectivity of the AR process and consequently the vanishing of the propagating AR (VAR process). In particular, the concept of the VAR process is a new physical process presented in this paper. If we would not accept the existence of this process, the total reflection independent of E is naively expected. Since finite transmission is possible in this angle region ($\theta_{c2} < |\theta| < \theta_{c1}$) above T_c , this total reflection would induce a sudden decrease of the conductance just below T_c for highly polarized ferromagnets junctions. As far as we know, no trends for such effect has been reported thus far. This fact may be the direct evidence for the existence of the VAR process. The VAR process is shown to have an important role on the Josephson current in superconductor/ferromagnet/superconductor junctions, because the evanescent wave carries a net Josephson current in this configuration³⁰. Note that the suppression mechanism of the AR process presented here is essentially different from that discussed in one-dimensional model where the contribution of the AR to the net current is simply governed by the ratio $k_{N,\downarrow}/k_{N,\uparrow}$ ⁸.

III. RESULTS

A. Effects of polarization

In this subsection, to reveal the influence of the polarization on the tunneling conductance spectra, we assume F/I/S junction by setting $\hat{U}_B = 0$ ($Z_{0,\uparrow} = Z_{0,\downarrow} \equiv Z_0$). At first, let us discuss several analytical results obtained from above formulation in order to check the validity of the formula. When $U = 0$, the ferromagnet reduces to a normal metal, and as expected $\sigma_q(E)$ reproduces the results of Ref.^{13,15}, and $\sigma_s(E)$ vanishes. For half-metallic ferromagnets ($U = E_{FN}$), the Fermi-surface for the down spins has shrunk to zero. In this case, the VAR process occurs for all θ_N . Under the condition of VAR, $\hat{\sigma}_q(E, \theta_N) = \hat{\sigma}_s(E, \theta_N)$ applies, which corresponds to the fact that the tunneling current is completely spin-polarized. Furthermore, the conductance spectra in the energy gap ($E < |\Delta_+|$, $E < |\Delta_-|$) become completely zero [$\sigma_q(E, \theta_N) = \sigma_s(E, \theta_N) = 0$]. In the tunneling limit ($H \rightarrow \infty$) and in the absence of VAR, $\hat{\sigma}_{q,\uparrow[\downarrow]}(E, \theta)$ gives the angle resolved surface DOS of an isolated superconductor. Then $\sigma_q(E)$ converges to the surface DOS weighted by the tunneling probability distribution¹⁵. At this limit, we can reproduce a well-known result that the ratio of the peak heights in the spin-resolved spectra directly reflect the polarization in the ferromagnet¹. On the other hand, $\sigma_{s,\uparrow[\downarrow]}(E, \theta)$ reduces to a function similar to the surface DOS, but where the divergence at the energy levels of the surface bound states is missing.

Next, the calculated results based on above formula are presented for $d_{x^2-y^2}$ -wave superconductors. In the following, we assume $E_{FN} = E_{FS}$. Figures 2 and 3 show the conductance spectra of charge current for the transparent limit ($Z_0 = 0$, $\beta = 0$) and high-barrier case ($Z_0 = 5$, $\beta = \pi/4$) as the function of exchange interaction $X (\equiv U/E_{FN})$. For $X = 0$, results in Ref.¹³ are reproduced. However, as X increases, the conductance inside the gap ($|E| < \Delta_0$) is largely reduced for both cases. Especially, the ZBCP disappears for the half-metallic ferromagnet case. Since the spin-polarization has such a drastic influence on the ZBCP, the height of ZBCP can be used in principle as a measurement of the magnitude of the spin polarization. Figure 4 shows the difference of the spin current and the charge current when $X=0.85$, $Z_0 = 5$ and $\beta = \pi/4$. It is clear that the ZBCP is not present for the spin current. This corresponds to the fact that the charge current components corresponding to the ZES are carried by condensed Cooper pairs in the superconductor, and therefore they do not contribute to the spin imbalance. As a result, the spin current becomes relatively insensitive to the orientation of the junctions. Figure 5 shows the conductance spectra for the spin current as the function of spin polarization ($Z_0 = 5$). It is clear that the spin current increases as X becomes larger. Note that $\sigma_s(E)$ is larger than unity around $E = \Delta_0$ when $X \approx 1$. This corresponds to the fact that the peak in the DOS has an influence even for the spin current.

Next, the net polarization $J_p(eV)$ is calculated for $d_{x^2-y^2}$ -wave superconductors as a function of the orientation (β) when $T = 0$. Four lines of Fig. 6 show the results for various values of barrier parameter when $eV = 2\Delta_0$. It is clear that the orientational effect is much smaller compared to the effect of Z_0 . In the same figure, results for s -wave superconductors ($\Delta_+ = \Delta_- = \Delta_0$ independent of θ_N) are also shown as closed dots. The large deviations of $d_{x^2-y^2}$ -wave from s -wave for small values of Z_0 are originated from the distribution of the pair amplitude in k -space. As the barrier parameter becomes larger, the spin injection

efficiency becomes to be insensitive to the symmetry of the pair potential.

B. Spin filtering effects and the ZBCP splitting

It has been experimentally verified that a ferromagnetic semiconductor used as the insulator in tunneling junctions works as a ferromagnetic barrier. Since the transmission probabilities for up and down spins are not equal, a spin-filtering effect is expected to be realized^{6,7}. Also it has been theoretically verified that a ferromagnetic insulator placed in the vicinity of superconductor induces a spin-splitting on the DOS of *s*-wave superconductors³¹. In the following, we will analyze the influence of the exchange interaction existing inside the insulator on the transport properties based on the formulation described in Sec. II.

Figure 7 shows the response of the conductance spectra $\sigma_q(E)$ on the exchange interaction in the insulator when $X = 0$. ZBCP splittings are obtained for finite exchange amplitude (U_B) cases. As U_B is increased and consequently as the difference between $Z_{0,\uparrow}$ and $Z_{0,\downarrow}$ becomes larger, the amplitude of the splitting becomes larger and the two peaks become broader and smaller. The peaks in the gap disappear when the difference between $Z_{0,\uparrow}$ and $Z_{0,\downarrow}$ becomes prominent. To see more clearly these trends, the spin-resolved conductance spectra $\sigma_{q,\uparrow[\downarrow]}(E)$ and $\sigma_s(E)$ for $Z_{0,\uparrow} = 2.5$ and $Z_{0,\downarrow} = 7.5$ are plotted in Fig. 8. The spectra for up [down] spins are shifted for lower [higher] energy level. Furthermore, $\sigma_s(E)$ becomes finite even though $X = 0$ in the ferromagnet. In order to check the effect of the polarization, Fig. 9 shows the response of the charge current as a function of polarization X for a fixed barrier parameter. The spin polarization in the ferromagnet induces the imbalance of the peak heights, thus the ratio of the splitted peak heights can be used as a criteria for the spin-polarization.

These results are interpreted as follows: i) The peaks corresponding to the up [down] spin components are shifted because of the energy gain (loss) during the tunneling process. ii) Since this energy gain (loss) has *k*-dependence³¹, the peak becomes broader comparing to the magnetic-field induced peak splitting (see below). iii) The amplitude of the peak splitting depends on the genuine barrier amplitude \hat{V}_0 as well as the exchange amplitude \hat{U}_B . For example, the splitted peaks merge into a single peak at the tunneling limit ($\hat{V}_0 \rightarrow \infty$) even if \hat{U}_B is kept constant. iv) The current corresponding to the ZBCP is carried by the Cooper pair in the superconductor as described in the previous subsection. This corresponds to the fact that the AR process is the second-lowest order tunneling process which requires both up and down spins tunneling. Hence, as $Z_{0,\downarrow}$ becomes larger and as the tunneling probability for down spins are suppressed, the conductance peaks and the AR process are rapidly reduced even if $Z_{0,\uparrow}$ is kept zero. v) The spin current is increased as \hat{U}_B is raised from zero even if X in the ferromagnet is kept at zero (unpolarized). This feature directly corresponds to the spin-filtering effect that the spin-selective tunneling occurs due to the presence of the exchange field in the insulator.

Next, various types of the ZBCP splitting expected for *d*-wave superconductors and their polarization effects are analyzed. Mainly two possibilities other than the ferromagnetic insulator effects have been proposed for the origins of the ZBCP splitting on high- T_c superconductor junctions. One is the Zeeman effect due to an applied magnetic field, and the other is the inducement of the BTRS states such as $d_{x^2-y^2}+is$ -wave. The conductance

spectra in an applied magnetic field is calculated from above formula by simply using the relation

$$\sigma_{q[s]}(E) = \sigma_{q[s],\uparrow}(E - \mu_B H) + \sigma_{q[s],\downarrow}(E + \mu_B H), \quad (17)$$

where $\mu_B H$ is the Zeeman energy. Calculated charge conductance spectra for $d_{x^2-y^2}$ -wave superconductor as a function of X are shown in Fig. 10. The amplitude of the splitting is linear to the applied field independent of the barrier heights. Moreover, since the energy shift induced by the magnetic field does not have k -dependence, the broadening of the peaks are not observed. The ratio of the splitted peak heights simply reflects the polarization in the ferromagnet, which is consistent with the results by Tedrow and Meservey¹. On the other hand, $\sigma_q(E)$ for $d_{x^2-y^2} + is$ -wave superconductor is calculated by setting $\Delta_{\pm} = \Delta_0 \cos 2(\theta_S \mp \beta) + i\Delta_s$. Calculated charge conductance spectra for various X values are shown in Fig. 11. The amplitude of the splitting is almost equivalent to the amplitude of the s -wave component. The shape of the spectrum without the polarization ($X = 0$) is quite similar to that shown in Fig. 10 ($X = 0$). As X becomes larger, the heights of the two peaks are reduced, which is consistent with that shown in Fig. 3. On the other hand, differently from Figs. 9 and 10, since the peak splitting is not induced by spin-dependent effects in this case, the polarization in the ferromagnets does not yield an imbalance in the peak heights. Thus the heights of the two peaks are reduced symmetrically.

The responses of the ZBCP on the variation of the polarization and the applied magnetic field are summarized as follows: i) the peak splitting due to the ferromagnetic insulator and the Zeeman effect are spin dependent. Therefore, the polarization in the ferromagnet induces the asymmetrical splitting of the ZBCP. ii) The amplitude of the peak splitting is linear to the applied field in the case of the Zeeman effect. However, it is non-linear in the cases of the ferromagnetic insulator effects⁷ and the BTRS states²³. In particular, the peak splittings are expected even in the absence of the applied field for these two cases. iii) The combination of the BTRS states and the Zeeman effect induces an additional peak splitting, that is, the ZBCP splits into four peaks. However, the combination of the Zeeman and the ferromagnetic insulator effects yields two peaks.

The experimental observations of the ZBCP splitting have been reported for normal metal / high- T_c superconductor junctions^{14,24,32,33}. It is a really interesting experiment to observe the same features by using ferromagnets/high- T_c superconductor junctions in order to distinguish the spin-dependent effects from the BTRS states inducement. Recently, Sawa, *et.al.* have detected an asymmetric magnetic field response in $\text{La}_{0.67}\text{Sr}_{0.33}\text{MnO}_3$ / $\text{YBa}_2\text{Cu}_3\text{O}_{7-\delta}$ junctions³⁴. The qualitative features on the magnetic field responses of their junctions are consistent with F/FI/S with $d_{x^2-y^2}$ -wave explained here. Detailed comparison between above formulas and their experiments is strongly expected.

Finally, a simple proposal is given for a possible device application utilizing the ferromagnetic insulator effects. The thickness of the insulator is the order of 1nm in usual tunneling junctions. Since the controlling of properties in such a thin layer requires high technology, as far as we know, not so many experimental trials have been accomplished thus far. However, as shown in this paper, a small change in the insulator property causes a drastic change on the transport properties. Therefore, the controlling of the barrier properties is one of the most promising methods to create new functional devices. For example, consider a F/FI/S junction with a $d_{x^2-y^2}$ -wave superconductor ($\beta = \pi/4$). The sharp ZBCP is drastically

modified as the difference between Z_{\uparrow} and Z_{\downarrow} becomes larger as shown in Fig. 7. This means that, for a fixed bias voltage, a large response in current is expected due to a small variation in the exchange interaction in the insulator. This response is applicable for the high-sensitive magnetization measurement of a thin insulating film by inserting the film into a junction as a tunneling barrier. If the exchange interaction is sensitive to the external field, this effect can be used as a magnetic sensor. Alternatively, if the magnetization of the insulator shows a hysteresis on the external field variation, a memory function can be realized. The current gain of the junction as a function of the external field is largely enhanced by using a superconductor /ferromagnetic insulator /superconductor junction with d -wave, because negative-conductance regions are expected just beside the ZBCP in this configuration^{35,36}. Differently from conventional superconducting memories based on a flux-quantum logic, a large-scale integration circuit may be possible based on the present principle.

IV. SUMMARY

In this paper, the conductance spectra for the charge and the spin currents under the influence of the exchange interaction have been calculated based on the scattering method. The influence of the spin polarization on the transport properties has been clarified. It is shown that the retro-reflectivity of the standard Andreev reflection process is broken in the presence of an exchange field and that the surface bound states due to superconducting pair potentials do not contribute to the spin current. Next, the ferromagnetic insulator including the spin-filtering effect are analyzed. It is shown that the spin-polarization gives asymmetric peak splitting. Moreover, various features in the splitting of ZBCP due to the ferromagnetic insulator, the Zeeman splitting, and the BTRS states effects are analyzed in detail. It is shown that the spin-polarized tunneling gives quite important information to identify the origin of the ZBCP splitting. By comparing the present analysis with experimental data, we expect that the mechanism of the peak splitting in high- T_c superconductors will be well identified. In the present model, we have neglected the effects of spin-orbit scattering³⁷ and the non-equilibrium properties of superconductors^{38,39}. Inclusion of these effects would be necessary for a complete theory. The formulation for triplet superconductors will be presented in another publication⁴⁰.

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FIGURES

FIG. 1. Schematic illustration of the elastic reflection of quasiparticles in the F/I/S junction. For all trajectories, momenta parallel to the interface are conserved. This means that the retro-reflection property of Andreev process is lost due to the exchange interaction. In the figure, the anisotropic pair potential of $d_{x^2-y^2}$ -wave symmetry is also shown.

FIG. 2. The normalized conductance spectra for the charge current $\sigma_q(E)$ as the function of $X \equiv U/E_{FN}$ with $\beta = 0$ and $Z_0 = 0$ (the transparent limit). As X becomes larger, the peak around zero-bias level is largely suppressed.

FIG. 3. The normalized conductance spectra for the charge current $\sigma_q(E)$ as the function of X with $\beta = \pi/4$ and $Z_0 = 5$. As X becomes larger, the height of the ZBCP is largely reduced.

FIG. 4. The comparison between the normalized conductance spectrum for the charge current $\sigma_q(E)$ and that for the spin current $\sigma_s(E)$ for $X = 0.7$, $\beta = \pi/4$ and $Z_0 = 5$. Since the ZBCP originates from the current carried by the surface bound states, the peak disappears for the spin current.

FIG. 5. The normalized conductance spectra for the spin current $\sigma_s(E)$ as the function of X with $\beta = 0$ and $Z_0 = 5$. As X becomes larger, the spin current is increased. Note that the peak at $E = \Delta_d$ is larger than unity when X is close to one.

FIG. 6. Orientational dependencies of $J_p(E)$ for $d_{x^2-y^2}$ -wave superconductors for $E = 2\Delta_0$ and $X = 0.7$ are plotted for various Z_0 values. Closed dots in the figures correspond to those for s -wave superconductors. The large deviations of $d_{x^2-y^2}$ -wave from s -wave for small values of Z_0 are originated from the distribution of the pair amplitude in k -space.

FIG. 7. The effects of ferromagnetic insulator on the charge current for $X = 0$ and $\beta = \pi/4$. When $Z_{0,\uparrow} = 5$ and $Z_{0,\downarrow} = 5$, a large ZBCP exists. The difference in $Z_{0,\uparrow}$ and $Z_{0,\downarrow}$ induces the peak splitting ($Z_{0,\uparrow} = 3$ and $Z_{0,\downarrow} = 7$). As the difference becomes larger, the ZBCP split into two peaks and the amplitude of the splitting becomes larger and the peaks become smaller and broader ($Z_{0,\uparrow} = 2$ and $Z_{0,\downarrow} = 8$). Finally, the peaks in the gap disappear as the difference becomes prominent ($Z_{0,\uparrow} = 0$ and $Z_{0,\downarrow} = 10$).

FIG. 8. Four types of normalized conductance for $Z_{0,\uparrow} = 2$, $Z_{0,\downarrow} = 8$, $X = 0$ and $\beta = \pi/4$. The normalized charge current $\sigma_q(E)$ has splitted peaks. The peak of the lower energy and higher energy are originated from the up spin component $\sigma_{q,\uparrow}(E)$ and down spin component $\sigma_{q,\downarrow}(E)$. These peaks do not appear in the spin current conductance $\sigma_s(E)$.

FIG. 9. The normalized conductance spectra $\sigma_q(E)$ as a function of X for $Z_{0,\uparrow} = 2.5$, $Z_{0,\downarrow} = 7.5$ and $\beta = \pi/4$. As X becomes larger, the higher energy peak becomes smaller. Thus, the polarization can be estimated from the ratio of two peak heights.

FIG. 10. The normalized conductance spectra $\sigma_q(E)$ in an applied magnetic field ($\mu_g H/\Delta_0 = 0.15$) as a function of X for $Z_{0,\uparrow} = 5$, $Z_{0,\downarrow} = 5$, and $\beta = \pi/4$. As X becomes larger, the peak with higher energy is largely reduced.

FIG. 11. The normalized conductance spectra $\sigma_q(E)$ for the BTRS states ($\Delta_s/\Delta_0 = 0.15$) as a function of X for $Z_{0,\uparrow} = 5$, $Z_{0,\downarrow} = 5$ and $\beta = \pi/4$. As X becomes larger, the heights of the two peaks are reduced symmetrically.





















